FORTY YEARS OF $\Lambda\Lambda$ HYPERNUCLEI*

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Recent experiments on production of $\Lambda\Lambda$ Hypernuclei have provoked renewed interest in extracting the $\Lambda\Lambda$ interaction from the few events identified since the inception of this field forty years ago. Few-body calculations relating to this issue are reviewed, particularly with respect to the possibility that $A = 4$ marks the onset of $\Lambda\Lambda$ binding to nuclei. The Nijmegen soft-core model potentials NSC97 qualitatively agree with the strength of the $\Lambda\Lambda$ interaction deduced from the newly determined binding energy of ${}_{\Lambda\Lambda}^{6}$He. Applying the extended NSC97 model to stranger nuclear systems suggests that $A = 6$ marks the onset of Ξ binding, with a particle stable ${}_{\Lambda\Xi}^{6}$He, and that strange hadronic matter is robustly bound.

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1. Introduction

Data on strangeness $S = -2$ hypernuclear systems is rather scarce, and no data exist for systems with higher strangeness content of hyperons (Y). Multistrange hadronic matter in finite systems and in bulk is predicted on general grounds to be stable, up to strangeness violating weak decays (see Ref. [1] for a recent review). Until 2001 only three candidates existed for $\Lambda\Lambda$ hypernuclei observed in emulsion experiments [2–4], the pioneering one of which, forty years old, is due to the leadership of the Warsaw distinguished experimental physicists Danysz and Pniewski [2]. The $\Lambda\Lambda$ binding energies deduced from these emulsion events indicated that the $\Lambda\Lambda$ interaction is strongly attractive in the ${}^1S_0$ channel [5–7], with a $\Lambda\Lambda$ pairing energy $\Delta B_{\Lambda\Lambda} \sim 4.5$ MeV, although it had been realized [8,9] that the binding energies of ${}_{\Lambda\Lambda}^{10}$Be [2] and ${}_{\Lambda\Lambda}^{6}$He [3] are inconsistent with each other. This outlook has undergone an important change following the very recent report from

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the KEK hybrid-emulsion experiment E373 on a well-established new candidate [10] for ${}_{\Lambda\Lambda}^{\;6}$He, with binding energy ($\Delta B_{\Lambda\Lambda} \sim 1$ MeV) substantially lower than that deduced from the older, dubious event [3]. Furthermore, there are also indications from the AGS experiment E906 for the production of light $\Lambda\Lambda$ hypernuclei [11], perhaps as light even as ${}_{\Lambda\Lambda}^{\;4}$H, in the (K^-, K^+) reaction on ^{9}Be. Further discussion of the latter experiment and its possible interpretations is provided by the companion talk of E.V. Hungerford.

The study of multistrange systems can provide stringent tests of microscopic models for the baryon-baryon (BB) interaction. The Nijmegen group has constructed a number of one-boson-exchange (OBE) models (reviewed in Ref. [12]) for the BB interaction using SU(3)-flavor symmetry to relate coupling constants and phenomenological short-distance hard or soft cores. In all of these rather different BB interaction models only 35 YN low-energy, generally imprecise data points serve the purpose of steering phenomenologically the extrapolation from the NN sector, which relies on thousands of data points, into the strange YN and YY sectors. It is therefore of utmost importance to confront these models with the new $\Lambda\Lambda$ hypernuclear data in order to provide meaningful constraints on the extrapolation to $S = -2$ and beyond.

2. $\Lambda\Lambda$ hypernuclei

In this section I will review topical work on some of the light $\Lambda\Lambda$ hypernuclear species connected to old and to new experiments. The anticipated existence of ${}_{\Lambda\Lambda}^{\;6}$He, now solidly established also experimentally [10], leads one to enquire where the onset of $\Lambda\Lambda$ binding occurs. It was argued long ago that the three-body $\Lambda\Lambda N$ system is unbound [13], and hence I will concentrate on the $A = 4, 5$ $\Lambda\Lambda$ hypernuclear systems. Among the few heavier species reported todate, ${}_{\Lambda\Lambda}^{10}$Be will be discussed briefly.

2.1. ${}_{\Lambda\Lambda}^{\;5}$H–${}_{\Lambda\Lambda}^{\;5}$He

Figure 1 demonstrates a nearly linear correlation between Faddeev-calculated values of $\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;6}\mathrm{He})$ and $\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;5}\mathrm{H},\ {}_{\Lambda\Lambda}^{\;5}\mathrm{He})$, using several $\Lambda\Lambda$ interactions including (the lowest-left point) $V_{\Lambda\Lambda} = 0$ [14]. Here

$$\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;A}Z) = B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;A}Z) - 2\bar{B}_{\Lambda}({}^{(A-1)}_{\quad\Lambda}Z)\,, \tag{1}$$

where $B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;A}Z)$ is the $\Lambda\Lambda$ binding energy of the hypernucleus ${}_{\Lambda\Lambda}^{\;A}Z$ and $\bar{B}_{\Lambda}\left({}^{(A-1)}_{\quad\Lambda}Z\right)$ is the $(2J+1)$-average of B_{Λ} values for the ${}^{(A-1)}_{\quad\Lambda}Z$ hypernuclear core levels. $\Delta B_{\Lambda\Lambda}$ increases monotonically with the strength of $V_{\Lambda\Lambda}$, starting in approximately zero as $V_{\Lambda\Lambda} \to 0$, which is a general feature of three-body

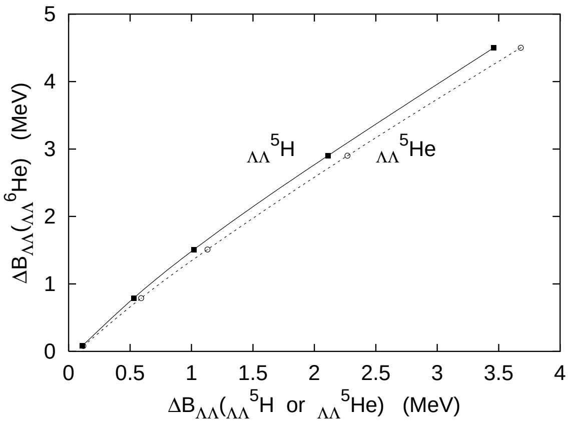

Fig. 1. s-wave Faddeev calculations [14] of $\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;6}\mathrm{He})$ *vs* $\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;5}\mathrm{H},\ {}_{\Lambda\Lambda}^{\;5}\mathrm{He})$.

models such as the $\alpha\Lambda\Lambda$, ${}^{3}\mathrm{H}\Lambda\Lambda$ and ${}^{3}\mathrm{He}\Lambda\Lambda$ models used in these s-wave Faddeev calculations [14], and also as shown below for $d\Lambda\Lambda$ s-wave Faddeev calculations [15]. The $I = 1/2$ ${}_{\Lambda\Lambda}^{\;5}\mathrm{H}$–${}_{\Lambda\Lambda}^{\;5}\mathrm{He}$ hypernuclei are then found to be particle stable for *all* the $\Lambda\Lambda$ attractive potentials here used. This conclusion holds also when the s-wave approximation is relaxed [16].

2.2. ${}_{\Lambda\Lambda}^{\;4}\mathrm{H}$

I start by discussing the first Faddeev–Yakubovsky four-body calculation of ${}_{\Lambda\Lambda}^{\;4}\mathrm{H}$ [15]. For two identical hyperons and two essentially identical nucleons (upon introducing isospin) as appropriate to a $\Lambda\Lambda pn$ model calculation of ${}_{\Lambda\Lambda}^{\;4}\mathrm{H}$, the 18 Faddeev–Yakubovsky components reduce to seven independent components satisfying coupled equations. Six rearrangement channels are involved in the s-wave calculation [15] for ${}_{\Lambda\Lambda}^{\;4}\mathrm{H}(1^{+})$:

$$(\Lambda NN)_{S=\frac{1}{2}} + \Lambda\ ,\ \ (\Lambda NN)_{S=\frac{3}{2}} + \Lambda\ ,\ \ (\Lambda\Lambda N)_{S=\frac{1}{2}} + N \qquad (2)$$

for 3+1 breakup clusters, and

$$(\Lambda\Lambda)_{S=0} + (NN)_{S=1}\ ,\quad (\Lambda N)_{S} + (\Lambda N)_{S'} \qquad (3)$$

with (S,S')=$(0,1)$+$(1,0)$ and $(1,1)$ for 2+2 breakup clusters.

Using $V_{\Lambda\Lambda}$ which reproduces $B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\;6}\mathrm{He})$, the four-body calculation converges well as function of the number N of the Faddeev–Yakubovsky basis functions allowed in, yet it yields no bound state for the $\Lambda\Lambda pn$ system, as demonstrated in Fig. 2 by the location of the '$\Lambda\Lambda pn$' curve *above* the hori-

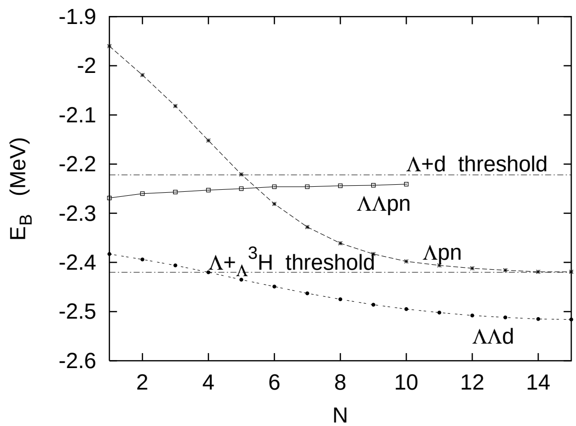

Fig. 2. s-wave Faddeev–Yakubovsky calculations [15] for Λpn, $\Lambda\Lambda d$ and $\Lambda\Lambda pn$.

zontal straight line marking the '$\Lambda+\ ^3_\Lambda$H threshold'[1]. In fact these Faddeev–Yakubovsky calculations exhibit little sensitivity to $V_{\Lambda\Lambda}$ over a wide range. Even for considerably stronger $\Lambda\Lambda$ interactions one gets a bound $^{\ 4}_{\Lambda\Lambda}$H only if the ΛN interaction is made considerably stronger, by as much as 40%. With four ΛN pairwise interactions out of a total of six, the strength of the ΛN interaction (about half of that for NN) plays a major role in the four-body $\Lambda\Lambda pn$ problem. However, fitting a Λd potential to the low-energy parameters of the s-wave Faddeev calculation for Λpn and solving the s-wave Faddeev equations for a $\Lambda\Lambda d$ model of $^{\ 4}_{\Lambda\Lambda}$H, this latter four-body system is calculated to yield a 1^+ bound state, as shown in the figure by the location of the asymptote of the '$\Lambda\Lambda d$' curve *below* the '$\Lambda+\ ^3_\Lambda$H threshold'. The onset of particle stability for $^{\ 4}_{\Lambda\Lambda}$H(1^+) requires then a minimum strength for $V_{\Lambda\Lambda}$ which is exceeded by the choice of $B_{\Lambda\Lambda}(^{\ 6}_{\Lambda\Lambda}$He) [10] as a normalizing datum. Disregarding spin it can be shown that, for essentially attractive $\Lambda\Lambda$ interactions and for a static nuclear core d, a two-body Λd bound state implies binding for the three-body $\Lambda\Lambda d$ system [18]. However, for a *non static* nuclear core d (made out of dynamically interacting proton and neutron), a Λd bound state does not necessarily imply binding for the $\Lambda\Lambda d$ system. It is questionable whether incorporating higher partial waves, and $\Lambda\Lambda - \Xi N$ coupling effects, will change this qualitative feature.

The above conclusions have been very recently challenged by Nemura *et al.* [19]. Fig. 3 demonstrates that within their stochastic variational calculation, which uses the NSC97f(FG) input of the Filikhin and Gal calcula-

[1] This threshold was obtained as the asymptote of the Λpn s-wave Faddeev calculation which uses model NSC97f [17] for the underlying ΛN interaction, yielding $B_\Lambda(^3_\Lambda\mathrm{H}(\frac{1}{2}^+)) = 0.19$ MeV. Using model NSC97e, with $B_\Lambda(^3_\Lambda\mathrm{H}) = 0.07$ MeV, does not alter the conclusions listed below.

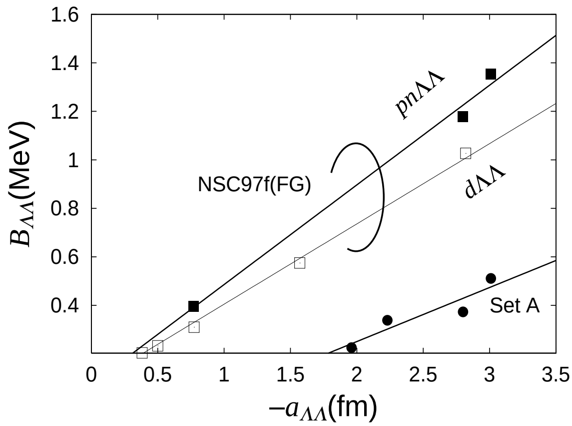

Fig. 3. $B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\ 4}\mathrm{H})$ as function of $a_{\Lambda\Lambda}$, calculated by Nemura *et al.* [19].

tion [15] for the various pairwise interactions, a '$pn\Lambda\Lambda$' model *always* yields more binding than a '$d\Lambda\Lambda$' model does. Particle stability for ${}_{\Lambda\Lambda}^{\ 4}\mathrm{H}(1^+)$ in this variational calculation requires a minimum strength for $V_{\Lambda\Lambda}$ which is exceeded by the choice of $B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\ 6}\mathrm{He})$ [10] as a normalizing datum. Yet, Nemura *et al.* argue that the ΛN interaction in the ${}^3S$ channel, when adjusted to the binding energy calculated for the $A = 4$ Λ hypernuclei, should be taken weaker than that used by Filikhin and Gal and that, when this constraint is implemented ('set A' in the figure), particle stability for ${}_{\Lambda\Lambda}^{\ 4}\mathrm{H}(1^+)$ requires a minimum strength for $V_{\Lambda\Lambda}$ which is not satisfied by the choice of $B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{\ 6}\mathrm{He})$ as a normalizing datum (equivalent to $-a_{\Lambda\Lambda} \sim 0.8$ fm [14]).

2.3. ${}_{\Lambda\Lambda}^{10}\mathrm{Be}$

For heavier $\Lambda\Lambda$ hypernuclei, the relationship between the three-body and four-body models is opposite to that found by Filikhin and Gal for ${}_{\Lambda\Lambda}^{\ 4}\mathrm{H}$: the $\Lambda\Lambda C_1C_2$ calculation provides *higher* binding than a properly defined $\Lambda\Lambda C$ calculation yields (with $C = C_1 + C_2$) due to the attraction induced by the ΛC_1–ΛC_2, $\Lambda\Lambda C_1$–C_2, C_1–$\Lambda\Lambda C_2$ four-body rearrangement channels that include bound states for which there is no room in the three-body $\Lambda\Lambda C$ model. The binding energy calculated within the four-body model increases then 'normally' with the strength of $V_{\Lambda\Lambda}$ [14]. This is demonstrated in Fig. 4 for ${}_{\Lambda\Lambda}^{10}\mathrm{Be}$ using several $\Lambda\Lambda$ interactions, including $V_{\Lambda\Lambda} = 0$ which corresponds to the lowest point on each one of the straight lines. The origin of the dashed axes corresponds to $\Delta B_{\Lambda\Lambda} = 0$. Within the 4-body $\alpha\alpha\Lambda\Lambda$ model, the fairly large value $\Delta B_{\Lambda\Lambda}({}_{\Lambda\Lambda}^{10}\mathrm{Be}) \sim 1.5$ MeV in the limit $V_{\Lambda\Lambda} \to 0$ is due to the special $\alpha\alpha$ cluster structure of the ${}^8\mathrm{Be}$ core. The correlation noted in the figure between ${}_{\Lambda\Lambda}^{10}\mathrm{Be}$ and ${}_{\Lambda\Lambda}^{\ 6}\mathrm{He}$ calculations, and the consistency

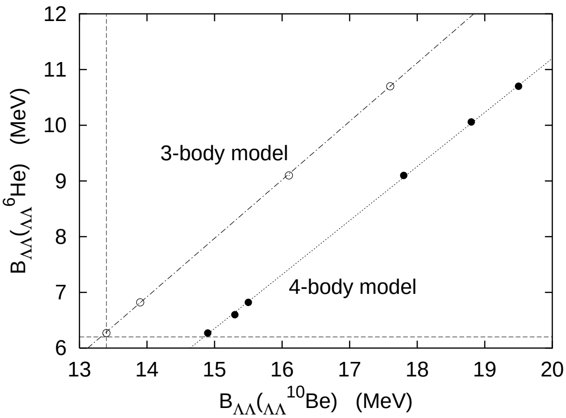

Fig. 4. s-wave Faddeev–Yakubovsky calculations [14] for $^{10}_{\Lambda\Lambda}\mathrm{Be}$: $^{8}\mathrm{Be}\Lambda\Lambda$ *versus* $\alpha\alpha\Lambda\Lambda$.

between various reports on their $B_{\Lambda\Lambda}$ values, are discussed by Filikhin and Gal [14, 20]. In particular, the two solid points next to the lowest one on the '4-body model' line in Fig. 4, corresponding to two versions of model NSC97 [17], are close to reproducing (the 'new') $B_{\Lambda\Lambda}(^{6}_{\Lambda\Lambda}\mathrm{He})$ but are short of reproducing (the 'old') $B_{\Lambda\Lambda}(^{10}_{\Lambda\Lambda}\mathrm{Be})$ by about 2.3 ± 0.4. This apparent discrepancy may be substantially reduced by accepting a $^{10}_{\Lambda\Lambda}\mathrm{Be}$ weak decay scheme that involves the 3 MeV excited $^{9}_{\Lambda}\mathrm{Be}$ doublet rather than the $^{9}_{\Lambda}\mathrm{Be}$ ground state [21]. This conclusion may also be inferred from the recent 4-body calculations by Hiyama *et al.* for $A = 7$–10 $\Lambda\Lambda$ hypernuclei [22].

3. The onset of Ξ stability

If model NSC97 [17] indeed provides a valid extrapolation from fits to NN and YN data, and noting the strongly attractive $^{1}S_{0}$ $\Lambda\Xi$ potentials in the extrapolation of this model to $S = -3, -4$ [23], it is natural to search for stability of $A = 6, S = -3$ systems obtained from $^{6}_{\Lambda\Lambda}\mathrm{He}$ upon replacing one of the Λ's by Ξ. Faddeev calculations [20] for the 0^{+} $I = 1/2$ ground-state $^{6}_{\Lambda\Xi}\mathrm{H}$ and $^{6}_{\Lambda\Xi}\mathrm{He}$, considered as $\alpha\Lambda\Xi^{-}$ and $\alpha\Lambda\Xi^{0}$ three-body systems respectively, indicate that $^{6}_{\Lambda\Xi}\mathrm{He}$ is particle-stable against Λ emission to $^{5}_{\Lambda\Lambda}\mathrm{He}$ for potentials simulating model NSC97, particularly versions e and f, whereas $^{6}_{\Lambda\Xi}\mathrm{H}$ is unstable since $M(\Xi^{-}) > M(\Xi^{0})$ by 6.5 MeV.[2] This is demonstrated in Fig. 5. Nevertheless, predicting particle stability for $^{6}_{\Lambda\Xi}\mathrm{He}$ is not independent of the assumptions made on the experimentally unexplored $\Xi\alpha$ interaction which was extrapolated from recent data on $^{12}\mathrm{C}$ [24]; hence this prediction cannot be considered conclusive.

[2] Recall that the $I = 1/2$ $^{5}_{\Lambda\Lambda}\mathrm{H}$–$^{5}_{\Lambda\Lambda}\mathrm{He}$ hypernuclei, within a $\Lambda\Lambda C$ Faddeev calculation, are particle stable even in the limit $V_{\Lambda\Lambda} \to 0$.

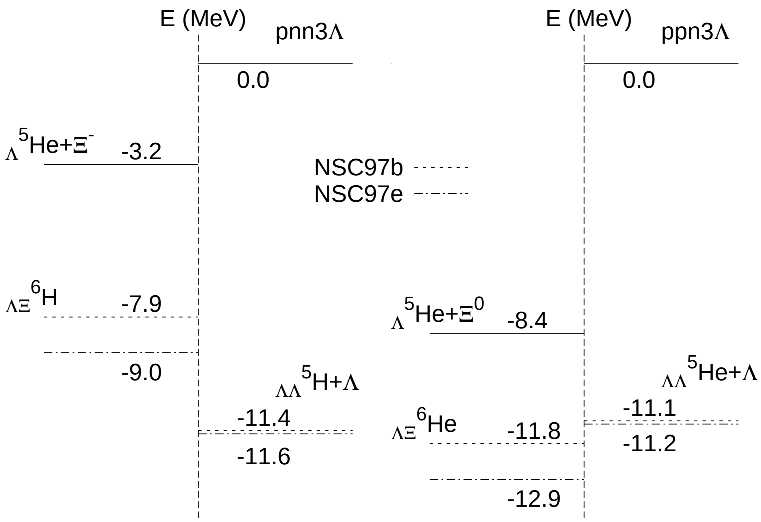

Fig. 5. s-wave Faddeev calculations [20] for $_{\Lambda\Xi}^{\;6}$H and $_{\Lambda\Xi}^{\;6}$He.

4. Strange hadronic matter

Bodmer [25], and more specifically Witten [26], suggested that strange quark matter, with roughly equal composition of u, d and s quarks, might provide the absolutely stable form of matter. Metastable strange quark matter has been studied by Chin and Kerman [27]. Jaffe and collaborators [28, 29] subsequently charted the various scenarios possible for the stability of strange quark matter, from absolute stability down to metastability due to weak decays. Finite strange quark systems, so called strangelets, have also been considered [28, 30].

Less advertised, perhaps, is the observation made by Schaffner *et al.* [31,32] that metastable strange systems with similar properties, *i.e.* a strangeness fraction $f_S = -S/A \approx 1$ and a charge fraction $f_Q = Z/A \approx 0$, might also exist in the hadronic basis at moderate values of density, between twice and three times nuclear matter density. These strange systems are made out of nucleons (N), lambda (Λ) and cascade (Ξ) hyperons. The metastability of these strange hadronic systems was established by extending relativistic mean field (RMF) calculations from ordinary nuclei ($f_S = 0$) to multi-strange nuclei with $f_S \neq 0$. Although the detailed pattern of metastability, as well as the actual values of the binding energy, depend specifically on the partly unknown hyperon potentials assumed in dense matter, the predicted phenomenon of metastability turned out to be robust in these calculations [33].

Recently, model NSC97 and its extension [23] were used to calculate within the RMF framework the minimum-energy equilibrium composition of bulk strange hadronic matter (SHM) made out of the SU(3) octet baryons N, Λ, Σ and Ξ, over the entire range of strangeness fraction $0 \leq f_S \leq 2$ [1]. The main result is that SHM is comfortably metastable in this model for any allowed value of $f_S > 0$. The $N\Lambda\Xi$ composition and the binding energy calculated for equilibrium configurations with $f_S \leq 1$ resemble those of model 2 in Refs. [31, 32]. The extension of model NSC97 [23] yields particularly attractive $\Xi\Xi$, $\Sigma\Sigma$ and $\Sigma\Xi$ interactions, but vanishingly weak $\Lambda\Lambda$ and $N\Xi$ interactions. Consequently, for $f_S \geq 1$, Σ's replace Λ's due to their exceptionally strong attraction to Σ and Ξ hyperons. As is shown below, a first-order phase transition occurs from $N\Lambda\Xi$ dominated matter for $f_S \leq 1$ to $N\Sigma\Xi$ dominated matter for $f_S \geq 1$, with binding energies per baryon reaching as much as 80 MeV.

A phase transition is visualized in Fig. 6 where the binding energy is drawn *versus* the baryon density for several representative fixed values of f_S. For $f_S = 0.8$, there is a global minimum at a baryon density of $\rho_B = 0.27$ fm^{-3}. A shallow local minimum is seen at larger baryon density at $\rho_B = 0.72$ fm^{-3}. Increasing the strangeness fraction to $f_S = 0.9$ lowers substantially the local minimum by about 20 MeV, whereas the global minimum barely changes. At $f_S = 1.0$ this trend is amplified and the relationship between the two minima is reversed, as the minimum at higher baryon density becomes energetically favored. The system will then undergo a transition from the low density state to the high density state. Due to the barrier between the two minima, it is a first-order phase transition from one minimum to the other.

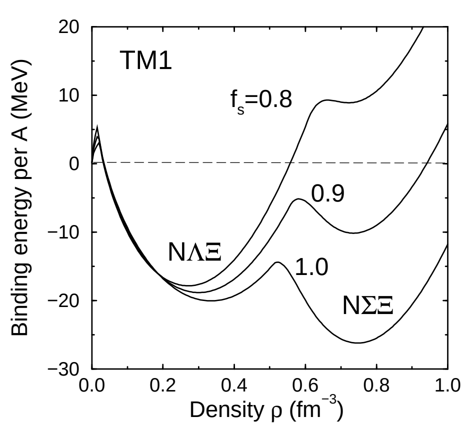

Fig. 6. Transition from $N\Lambda\Xi$ to $N\Sigma\Xi$ matter upon increasing the strangeness fraction [1].

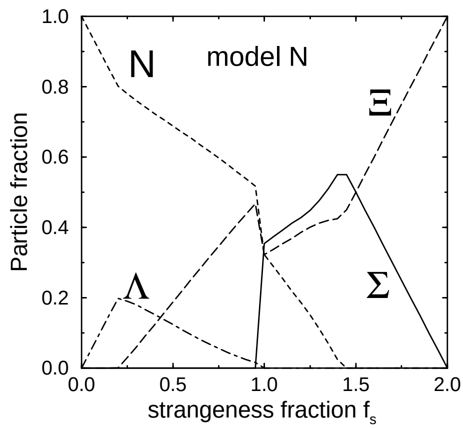

Fig. 7. Composition of strange hadronic matter *vs* the strangeness fraction [1].

Fig. 7 demonstrates explicitly that the phase transition involves transformation from $N\Lambda\Xi$ dominated matter to $N\Sigma\Xi$ dominated matter, by showing the calculated composition of SHM for this model (denoted N) as function of the strangeness fraction f_S. The particle fractions for each baryon species change as function of f_S. At $f_S = 0$, one has pure nuclear matter, whereas at $f_S = 2$ one has pure Ξ matter. In between, matter is composed of baryons as dictated by chemical equilibrium. A change in the particle fraction may occur quite drastically when new particles appear, or existing ones disappear in the medium. A sudden change in the composition is seen in Fig. 7 for $f_S = 0.2$ when Ξ's emerge in the medium, or at $f_S = 1.45$ when nucleons disappear. The situation at $f_S = 0.95$ is a special one, as Σ's appear in the medium, marking the first-order phase transition observed in the previous figure. The baryon composition alters completely at that point, from $N\Xi$ baryons plus a rapidly vanishing fraction of Λ's into $\Sigma\Xi$ hyperons plus a decreasing fraction of nucleons. At the very deep minimum of the binding energy curve (Fig. 3 of Ref. [1]) the matter is composed mainly of Σ's and Ξ's with a very small admixture of nucleons.

5. Conclusion

I have presented Faddeev calculations for ${}^{5}_{\Lambda\Lambda}$H–${}^{5}_{\Lambda\Lambda}$He and ${}^{6}_{\Lambda\Lambda}$He, and first ever four-body Faddeev–Yakubovsky calculations for ${}^{4}_{\Lambda\Lambda}$H and ${}^{10}_{\Lambda\Lambda}$Be, using NN and ΛN interaction potentials within the ${}^{4}_{\Lambda\Lambda}$H calculation that fit the available data on the relevant subsystems, including the binding energy of ${}^{3}_{\Lambda}$H. No ${}^{4}_{\Lambda\Lambda}$H bound state was obtained for a wide range of $\Lambda\Lambda$ interactions, including that corresponding to $B_{\Lambda\Lambda}({}^{6}_{\Lambda\Lambda}\mathrm{He})$. This non binding is due to the relatively weak ΛN interaction, in stark contrast to the results of a 'reasonable' three-body $\Lambda\Lambda d$ Faddeev calculation. More theoretical work, particularly on the effects of including explicitly $\Lambda\Lambda$–ΞN–$\Sigma\Sigma$ channel couplings, is called for. Preliminary estimates for such effects within the NSC97 model, or its simulation, have been recently made [16, 34–36]. Further experimental work is needed to decide whether or not the events reported in the AGS experiment E906 [11] correspond to ${}^{4}_{\Lambda\Lambda}\mathrm{H}$, particularly in view of subsequent conflicting theoretical analyses [37, 38].

Accepting the predictive power of model NSC97 which qualitatively reproduces $B_{\Lambda\Lambda}({}^{6}_{\Lambda\Lambda}\mathrm{He})$, Faddeev calculations suggest that ${}^{6}_{\Lambda\Xi}\mathrm{He}$ may be the lightest particle-stable $S = -3$ hypernucleus, and the lightest and least strange particle-stable hypernucleus in which a Ξ hyperon is bound. Unfortunately, the direct production of $\Lambda\Xi$ hypernuclei is beyond present experimental capabilities, requiring the use of Ω^- initiated reactions.

Finally, I have focused on the consequences of using model NSC97 for the binding and composition of strange hadronic matter. Strange hadronic matter is comfortably stable, up to weak decays, over a wide range of baryon-baryon interaction models, including model NSC97 here chosen because it nearly successfully extrapolates from the $S = 0, -1$ sectors in which it was constructed into the $S = -2$ sector. The phase transition considered in this review has been recently discussed by the Frankfurt group [39] in the context of phase transition to hyperon matter in neutron stars. Unfortunately, it will take lots of imagination to devise experimentally a way to determine how attractive those $\Lambda\Xi$, $\Xi\Xi$, $\Xi\Sigma$, $\Sigma\Sigma$ interactions are, which are so crucial for the results exhibited in this review.

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